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SINGLE SPIN ASYMMETRIES IN INCLUSIVE HADRON PRODUCTION [*]

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ABSTRACT

The general formalism to describe single spin asymmetries in hadron-hadron high energy and large p_T inclusive production within the QCD-improved parton model and assuming the factorization theorem to hold for higher twist contributions is discussed. Non zero values of the single spin asymmetries originate from and reveal non perturbative universal properties of quarks: the quark distribution and fragmentation analysing powers, which need not be zero, provided the quark intrinsic motion is taken into account. A simple model is constructed which reproduces the main features of the data on the single spin asymmetries observed in inclusive pion production in pp collisions.

According to the QCD factorization theorem the differential cross-section for the hard scattering of a polarized hadron A with spin S_A off another polarized hadron B with spin S_B, resulting in the inclusive production of a hadron C with energy E_C and three-momentum $\boldsymbol{p}_C$, $A, S_A + B, S_B \to C + X$, can be written as [1,2,3]

$$\frac{E_C \, d\sigma^{A,S_A+B,S_B\to C+X}}{d^3\boldsymbol{p}_C} = \sum_{a,b,c,d} \; \sum_{\lambda_a,\lambda'_a;\lambda_b,\lambda'_b;\lambda_c,\lambda'_c;\lambda_d;\lambda_C} \int \frac{dx_a \, dx_b}{\pi z} \frac{1}{16\pi\hat{s}^2} \tag{1}$$
$$\times \; \rho^{a/A,S_A}_{\lambda_a,\lambda'_a} \, f_{a/A}(x_a) \, \rho^{b/B,S_B}_{\lambda_b,\lambda'_b} \, f_{b/B}(x_b) \, \hat{M}_{\lambda_c,\lambda_d;\lambda_a,\lambda_b} \, \hat{M}^*_{\lambda'_c,\lambda_d;\lambda'_a,\lambda'_b} \, D^{\lambda_c,\lambda'_c}_{\lambda_C,\lambda_C}(z) \, ,$$

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where $f_{i/I}(x_i)$ is the number density of partons i with momentum fraction x_i inside hadron I ($i = a, b$; $I = A, B$) and $\rho^{i/I,S_I}_{\lambda_i,\lambda'_i}(x_i)$ is the helicity density matrix of parton i inside the polarized hadron I. The $\hat{M}_{\lambda_c,\lambda_d;\lambda_a,\lambda_b}$'s are the helicity amplitudes for the elementary process $ab \to cd$; if one wishes to consider higher order (in α_s) contributions also elementary processes involving more partons should be included. $D^{\lambda_c,\lambda'_c}_{\lambda_C,\lambda'_C}(z)$ is the product of *fragmentation amplitudes*

$$D^{\lambda_c,\lambda'_c}_{\lambda_C,\lambda'_C} = \sum\!\!\!\!\!\!\!\int_{X,\lambda_X} \mathcal{D}_{\lambda_X,\lambda_C;\lambda_c}\,\mathcal{D}^*_{\lambda_X,\lambda'_C;\lambda'_c}\,, \tag{2}$$

where the $\sum\!\!\!\!\!\!\int_{X,\lambda_X}$ stays for a spin sum and phase space integration of the undetected particles, considered as a system X. The usual unpolarized fragmentation function $D_{C/c}(z)$, *i.e.* the density number of hadrons C resulting from the fragmentation of an unpolarized parton c and carrying a fraction z of its momentum is given by

$$D_{C/c}(z) = \frac{1}{2}\sum_{\lambda_c,\lambda_C} D^{\lambda_c,\lambda_c}_{\lambda_C,\lambda_C}(z)\,. \tag{3}$$

For simplicity of notations we have not shown in Eq. (1) the Q^2 scale dependences in f and D; the variable z is related to x_a and x_b by the usual imposition of energy momentum conservation in the elementary $2 \to 2$ process [4], $z = -(x_a t + x_b u)/x_a x_b s$ where s, t, u are the Mandelstam variables for the overall process $AB \to CX$ whereas $\hat{s}, \hat{t}, \hat{u}$ are those for the elementary process $ab \to cd$. The corresponding amplitudes are normalized so that

$$\frac{d\hat{\sigma}}{d\hat{t}} = \frac{1}{16\pi\hat{s}^2}\frac{1}{4}\sum_{\lambda_a,\lambda_b,\lambda_c,\lambda_d} |\hat{M}_{\lambda_c,\lambda_d;\lambda_a,\lambda_b}|^2\,. \tag{4}$$

Eq. (1) holds at leading twist and large p_T values of the produced pion; the intrinsic $\boldsymbol{k}_\perp$ of the partons have been integrated over and collinear configurations dominate both the distribution functions and the fragmentation processes; one can then see that, in this case, there cannot be any single spin asymmetry.

Suppose, in fact, that hadron B is not polarized so that Eq. (1) reads

$$\begin{aligned}\frac{E_C\, d\sigma^{A,S_A+B\to C+X}}{d^3\boldsymbol{p}_C} \;&=\; \sum_{a,b,c,d}\;\sum_{\lambda_a,\lambda'_a;\lambda_b;\lambda_c,\lambda'_c;\lambda_d;\lambda_C} \frac{1}{2}\int \frac{dx_a\,dx_b}{\pi z}\frac{1}{16\pi\hat{s}^2} \\ &\times\; \rho^{a/A,S_A}_{\lambda_a,\lambda'_a}\, f_{a/A}(x_a)\, f_{b/B}(x_b)\, \hat{M}_{\lambda_c,\lambda_d;\lambda_a,\lambda_b}\, \hat{M}^*_{\lambda'_c,\lambda_d;\lambda'_a,\lambda_b}\, D^{\lambda_c,\lambda'_c}_{\lambda_C,\lambda_C}(z)\,.\end{aligned} \tag{5}$$

Then total angular momentum conservation in the (forward) fragmentation process [see Eq. (2)] implies $\lambda_c = \lambda'_c$; this, in turns, together with helicity conservation in the elementary processes, implies $\lambda_a = \lambda'_a$. If we further notice that, by parity invariance, $\sum_{\lambda_C} D^{\lambda_c,\lambda_c}_{\lambda_C,\lambda_C} = D_{C/c}$ independently of λ_c and that

$$\sum_{\lambda_b,\lambda_c,\lambda_d} |\hat{M}_{\lambda_c,\lambda_d;\lambda_a,\lambda_b}|^2 = 2\left(16\pi\hat{s}^2\right)\frac{d\hat{\sigma}}{d\hat{t}}\,, \tag{6}$$

we remain with $\sum_{\lambda_a} \rho^{a/A,S_A}_{\lambda_a,\lambda_a} = 1$. Moreover, in the absence of intrinsic $\boldsymbol{k}_\perp$ and initial state interactions, the parton density numbers $f_{a/A}(x_a)$ cannot depend on the hadron spin and any spin dependence disappears from Eq. (1), so that

$$\begin{aligned}\frac{E_C\, d\sigma^{A,S_A+B\to C+X}}{d^3\boldsymbol{p}_C} &= \sum_{a,b,c,d} \int \frac{dx_a\, dx_b}{\pi z}\, f_{a/A}(x_a)\, f_{b/B}(x_b) \\ &\times \frac{d\hat{\sigma}^{a+b\to c+d}}{d\hat{t}}\, D_{C/c}(z) \;=\; \frac{E_C\, d\sigma^{unp}}{d^3\boldsymbol{p}_C}\,. \end{aligned} \tag{7}$$

This implies that the single spin asymmetry,

$$\begin{aligned} A_N &\equiv \frac{d\sigma^{A^\uparrow B\to C\,X}(\boldsymbol{p}_T) - d\sigma^{A^\downarrow B\to C\,X}(\boldsymbol{p}_T)}{d\sigma^{A^\uparrow B\to C\,X}(\boldsymbol{p}_T) + d\sigma^{A^\downarrow B\to C\,X}(\boldsymbol{p}_T)} \\ &= \frac{d\sigma^{A^\uparrow B\to C\,X}(\boldsymbol{p}_T) - d\sigma^{A^\uparrow B\to C\,X}(-\boldsymbol{p}_T)}{d\sigma^{A^\uparrow B\to C\,X}(\boldsymbol{p}_T) + d\sigma^{A^\uparrow B\to C\,X}(-\boldsymbol{p}_T)}\,, \end{aligned} \tag{8}$$

is zero. $\uparrow$ ($\downarrow$) refers to hadron A spin orientation up (down) with respect to the production plane; any other spin orientation with no up or down component would result in a vanishing asymmetry due to parity invariance.

Eq. (1) can be generalized with the inclusion of intrinsic $\boldsymbol{k}_\perp$ [2] and this can avoid the above conclusion; for example [2], the observation of a non zero $\boldsymbol{k}_\perp$ of a final particle C with respect to the axis of the jet generated by parton c does not imply any more $\lambda_c = \lambda'_c$ and allows a non zero value of the asymmetry through the *quark fragmentation analysing power*

$$\begin{aligned} A_{C/q} &= \frac{D_{C/q,s_q}(z,\boldsymbol{k}_\perp) - D_{C/q,-s_q}(z,\boldsymbol{k}_\perp)}{D_{C/q,s_q}(z,\boldsymbol{k}_\perp) + D_{C/q,-s_q}(z,\boldsymbol{k}_\perp)} \\ &= \frac{D_{C/q,s_q}(z,\boldsymbol{k}_\perp) - D_{C/q,s_q}(z,-\boldsymbol{k}_\perp)}{D_{C/q,s_q}(z,\boldsymbol{k}_\perp) + D_{C/q,s_q}(z,-\boldsymbol{k}_\perp)}\,. \end{aligned} \tag{9}$$

A non zero value of the above quantity – the analysing power or single spin asymmetry for the quark fragmentation process – is allowed by parity invariance for quark spin orientations perpendicular to the q-C plane and is allowed by time-reversal invariance due to the soft interactions of the fragmenting quark with external fields.

This idea was exploited in Ref. 5 to explain the single spin asymmetries observed in $p^\uparrow p \to \pi X$ processes [6] and where, essentially, it is assumed that parton c is produced in the forward direction and the final hadron p_T is due to its transverse $k_\perp$ inside the jet. One cannot expect such a model to work at large p_T, as the results clearly show.

Another possible $\boldsymbol{k}_\perp$ effect, suggested by Sivers [7,8], may originate in the distribution functions. Similarly to the Collins effect in the fragmentation process a dependence on the hadron spin may remain in the $\boldsymbol{k}_\perp$ dependent quark distribution function $\tilde{f}_{q/A^\uparrow}(x,\boldsymbol{k}_\perp)$, so that the difference

$$\begin{aligned} \Delta^N \tilde{f}_{q/A^\uparrow}(x,\boldsymbol{k}_\perp) &\equiv \tilde{f}_{q/A^\uparrow}(x,\boldsymbol{k}_\perp) - \tilde{f}_{q/A^\downarrow}(x,\boldsymbol{k}_\perp) \\ &= \tilde{f}_{q/A^\uparrow}(x,\boldsymbol{k}_\perp) - \tilde{f}_{q/A^\uparrow}(x,-\boldsymbol{k}_\perp) \end{aligned} \tag{10}$$

can be different from zero. $\tilde{f}_{q/A^\uparrow}(x, \boldsymbol{k}_\perp)$ is the number density of partons a with momentum fraction x_a and intrinsic transverse momentum $\boldsymbol{k}_\perp$ in a transversely polarized hadron A.

Taking intrinsic $\boldsymbol{k}_\perp$ into account Eq. (7) becomes

$$\begin{aligned} \frac{E_C \, d\sigma^{A^\uparrow B \to CX}}{d^3\boldsymbol{p}_C} &= \sum_{a,b,c,d} \int \frac{d^2\boldsymbol{k}_\perp dx_a \, dx_b}{\pi z} \\ &\quad \times \, \tilde{f}_{a/A^\uparrow}(x_a, \boldsymbol{k}_\perp) \, f_{b/B}(x_b) \, \frac{d\hat{\sigma}}{d\hat{t}}(\boldsymbol{k}_\perp) \, D_{C/c}(z) \, , \end{aligned} \tag{11}$$

so that, adopting obvious shorter notations,

$$\begin{aligned} d\sigma^{A^\uparrow B \to CX} - d\sigma^{A^\downarrow B \to CX} &= \sum_{a,b,c,d} \int \frac{d^2\boldsymbol{k}_\perp dx_a \, dx_b}{\pi z} \qquad (12) \\ &\times \left[\tilde{f}_{a/A^\uparrow}(x_a, \boldsymbol{k}_\perp) - \tilde{f}_{a/A^\uparrow}(x_a, -\boldsymbol{k}_\perp) \right] \, f_{b/B}(x_b) \, d\hat{\sigma}(\boldsymbol{k}_\perp) \, D_{C/c}(z) \, . \end{aligned}$$

The above new quantity (10), divided by twice the unpolarized $\boldsymbol{k}_\perp$ dependent distribution function $2\tilde{f}_{q/A}(x_a, \boldsymbol{k}_\perp) = \tilde{f}_{q/A^\uparrow}(x_a, \boldsymbol{k}_\perp) + \tilde{f}_{q/A^\downarrow}(x_a, \boldsymbol{k}_\perp)$, can be regarded as a single spin asymmetry or analysing power for the $A^\uparrow \to a + X$ process; similarly to the quark fragmentation analysing power, Eq. (9), it may be different from zero for hadron spin orientations perpendicular to the A-q plane and taking into account initial state interactions between the two colliding hadrons. This quantity plays, for single spin asymmetries, the same role plaid by distribution functions in unpolarized processes. If we define the polarized number densities in terms of *distribution amplitudes* as

$$\tilde{f}_{a,\lambda_a/A^\uparrow}(x_a, \boldsymbol{k}_{\perp a}) = \sum\!\!\!\!\!\!\!\int_{X_A, \lambda_{X_A}} |\mathcal{G}^{a/A}_{\lambda_{X_A}, \lambda_a; \uparrow}(x_a, \boldsymbol{k}_{\perp a})|^2 \, , \tag{13}$$

then we have, in the helicity basis,

$$\begin{aligned} \Delta^N \tilde{f}_{a/A^\uparrow}(x_a, \boldsymbol{k}_{\perp a}) &= \sum\!\!\!\!\!\!\!\int_{X_A, \lambda_{X_A}} \sum_{\lambda_a} 2 \, \mathrm{Im} \left[\mathcal{G}^{a/A}_{\lambda_{X_A}, \lambda_a; +}(x_a, \boldsymbol{k}_{\perp a}) \, \mathcal{G}^{a/A \, *}_{\lambda_{X_A}, \lambda_a; -}(x_a, \boldsymbol{k}_{\perp a}) \right] \\ &\equiv 2 \, I^{a/A}_{+-}(x_a, \boldsymbol{k}_{\perp a}) \, , \end{aligned} \tag{14}$$

which shows the non diagonal nature, in the helicity indices, of $I^{a/A}_{+-}(x_a, \boldsymbol{k}_{\perp a})$.

Notice that, due to the odd $\boldsymbol{k}_\perp$ dependence of $I^{a/A}_{+-}$, one has

$$\begin{aligned} &\int d^2\boldsymbol{k}_\perp \, I^{a/A}_{+-}(\boldsymbol{k}_\perp) \, \frac{d\hat{\sigma}}{d\hat{t}}(\boldsymbol{k}_\perp) \\ &= \int_{(\boldsymbol{k}_\perp)_x > 0} d^2\boldsymbol{k}_\perp \, I^{a/A}_{+-}(\boldsymbol{k}_\perp) \, \left[\frac{d\hat{\sigma}}{d\hat{t}}(\boldsymbol{k}_\perp) - \frac{d\hat{\sigma}}{d\hat{t}}(-\boldsymbol{k}_\perp) \right] \, , \end{aligned} \tag{15}$$

where $d\hat{\sigma}/d\hat{t}(\boldsymbol{k}_\perp) - d\hat{\sigma}/d\hat{t}(-\boldsymbol{k}_\perp) = \mathcal{O}(k_\perp/p_T)$, which clearly shows that this is a higher twist effect.

A simple phenomenological model was developed in Ref. 9 for the $p^\uparrow p \to \pi X$ process by assuming that the dependence of $I_{+-}^{a/p}$ on $\boldsymbol{k}_{\perp a}$ is sharply peaked around an average value $k_{\perp a}^0$:

$$I_{+-}^{a/p}(x_a, k_x) = \frac{k_x}{M}\, \delta\left(|k_x| - k_{\perp a}^0\right) N_a\, x_a^{\alpha_a}(1-x_a)^{\beta_a} \tag{16}$$

with, from a fit to the data [10],

$$\frac{k_{\perp a}^0}{M} = 0.47\, x^{0.68}\,(1-x)^{0.48}\,, \tag{17}$$

where M is a hadronic mass scale, $M \simeq 1$ GeV/c.

Inserting Eqs. (12), (15) and (16) into Eq. (8) yields

$$A_N = \frac{1}{M}\, \frac{\sum N_a \int dx_a dx_b\, k_{\perp a}^0\, x_a^{\alpha_a}(1-x_a)^{\beta_a}\, f_{b/p}(x_b)\,[d\hat{\sigma}(k_{\perp a}^0) - d\hat{\sigma}(-k_{\perp a}^0)]\, D_{\pi/c}(z)/z}{\sum \int dx_a dx_b\; f_{a/p}(x_a)\, f_{b/p}(x_b) d\hat{\sigma}\, D_{\pi/c}(z)/z}\,. \tag{18}$$

In order to give numerical estimates of the asymmetry (18) we have taken $f_{q,\bar{q},g/p}$ from Ref. 11, $D_{\pi/q,\bar{q}}$ from Ref. 12 and $D_{\pi/g}$ from Ref. 13. Given the very limited p_T range of the data we have neglected the QCD Q^2 dependence of the distribution and fragmentation functions. We have only considered contributions from u and d quarks inside the polarized proton, which certainly dominate at large x_F values, that is $a = u, d$ in the numerator of Eq. (18). Instead, we have considered all possible constituents in the unpolarized protons, with $k_\perp = 0$, and all possible constituent fragmentation functions.

By using Eq. (17) into Eq. (18), together with the unpolarized $f_{a,b/p}$ and $D_{\pi/c}$ functions, we remain with an expression of A_N still dependent on a set of 6 free parameters, namely N_a, α_a and β_a ($a = u, d$), defined in Eq. (16). We have obtained a best fit to the data [6], shown in Fig. 1, with the following values of the parameters:

$$\begin{array}{cccc} & N_a & \alpha_a & \beta_a \\ u & 5.19 & 2.10 & 3.67 \\ d & -2.29 & 1.43 & 4.22 \end{array} \tag{19}$$

As the experimental data [6] cover a p_T range between 0.7 and 2.0 GeV/c we have computed A_N at a fixed value $p_T = 1.5$ GeV/c.

Notice that the above values (19) are very reasonable indeed; actually, apart from an overall normalization constant, they might even have been approximately

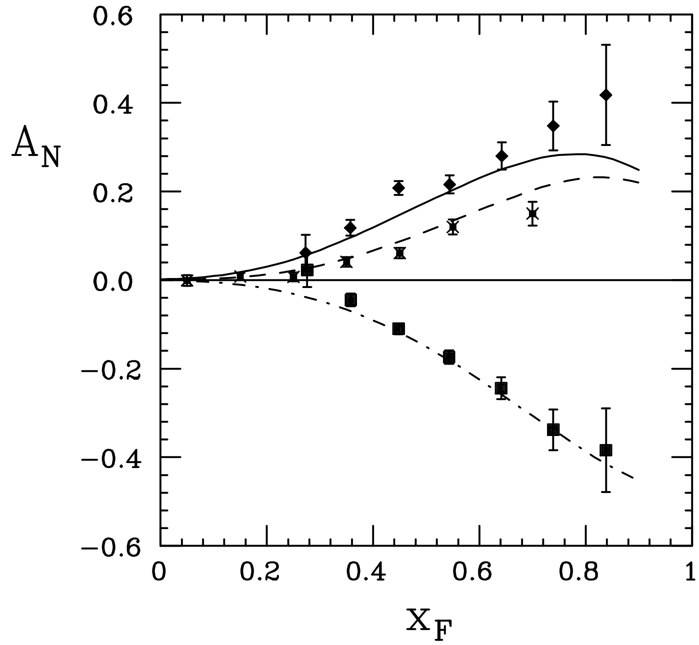

Fig. 1. Fit of the data on A_N [6], with the parameters given in Eq. (19); the upper, middle and lower sets of data and curves refer respectively to π^+, π^0 and π^-.

guessed. The exponents $\alpha_{u,d}$ and $\beta_{u,d}$ are not far from the very naïve values one can obtain by assuming, as somehow suggested by Eqs. (13) and (14), that $I_{+-}^{a/p}(x_a) \sim \sqrt{f_{a,+/p,+}(x_a) f_{a,-/p,+}(x_a)}$, where $f_{a,+(-)/p,+}(x_a)$ denotes, as usual, the number density of quarks with the same (opposite) helicity as the parent proton. Also the relative sign and strength of the normalization constants N_u and N_d turn out not to be surprising if one assumes that there might be a correlation between the number of quarks at a fixed value of $\boldsymbol{k}_{\perp a}$ and their polarization: remember that, according to $SU(6)$, inside a proton polarized along the $\hat{\boldsymbol{y}}$ direction, $P_y = 1$, one has for valence quarks $P_y^u = 2/3$ and $P_y^d = -1/3$.

It might appear surprising to have approximately opposite values for the π^+ and π^- asymmetries, as the data indicate, and a large positive value for the π^0; one might rather expect $A_N \simeq 0$ for a π^0. However, this can easily be understood from Eq. (18) which we simply rewrite, for a pion π^i, as $A_N^i = \mathcal{N}^i/\mathcal{D}^i$, if one remembers that from isospin symmetry one has:

$$D_{\pi^0/c} = \frac{1}{2}\left(D_{\pi^+/c} + D_{\pi^-/c}\right) . \tag{20}$$

Eq. (18) shows that the relation (20) also holds for $\mathcal{N}^0$ and $\mathcal{D}^0$, so that

$$A_N^0 = \frac{\mathcal{N}^+ + \mathcal{N}^-}{\mathcal{D}^+ + \mathcal{D}^-} = A_N^+ \; \frac{1 + \dfrac{A_N^-}{A_N^+}\dfrac{\mathcal{D}^-}{\mathcal{D}^+}}{1 + \dfrac{\mathcal{D}^-}{\mathcal{D}^+}} \; . \tag{21}$$

It is then clear that $A_N^- \simeq -A_N^+$ implies $A_N^0 \simeq 0$ only if $\mathcal{D}^- \simeq \mathcal{D}^+$, *i.e.* if the unpolarized cross-sections for the production of a π^- and a π^+ are approximately equal. This is true only at $x_F \simeq 0$. At large x_F the minimum value of x_a kinematically allowed increases and the dominant contribution to the production of π^+ and π^- comes respectively from $f_{u/p}(x_a)$ and $f_{d/p}(x_a)$ [see the denominator of Eq. (18)]. It is known that $f_{d/p}(x_a)/f_{u/p}(x_a) \to 0$ when $x_a \to 1$; this implies that $\mathcal{D}^-/\mathcal{D}^+$ decreases with increasing x_F, so that *at large* x_F we have $\mathcal{D}^-/\mathcal{D}^+ \ll 1$ and $A_N^0 \simeq A_N^+(1-2\mathcal{D}^-/\mathcal{D}^+) \simeq A_N^+$. Such a trend emerges both from the experimental data and our computations.

We have discussed how to compute single spin asymmetries in large p_T inclusive production within the framework of the factorization theorem and perturbative QCD, showing that single spin effects of order $k_\perp/p_T$ can be related to non perturbative intrinsic properties of quark fragmentations and distributions, *i.e.* the *quark fragmentation analysing power* and the *quark distribution analysing power*; the former was first suggested by Collins [2] and the latter by Sivers [7,8]. We have shown a possible description of the single spin asymmetries observed in $p^\uparrow p \to \pi X$ via the intrinsic $\boldsymbol{k}_\perp$ effects in the quark distribution functions.

A definite test and a better evaluation of these non perturbative properties requires further and more refined applications of the same idea and more theoretical work. In particular one might consider the following processes:

– $\bar{p}^\uparrow p \to \pi X$; a straightforward application of the model previously described, with the same set of parameters, gives results comparable with some existing data [14].

– $p^\uparrow p \to \gamma X$; in such a case there cannot be any fragmentation effect and one should be able to single out the importance of the Sivers effect and the quark distribution analysing power.

– $\ell p^\uparrow \to hX$; in such a case, the inclusive production of a hadron in DIS, there should be no effect from the quark distribution analysing power, as any initial state interaction would be of higher electromagnetic order and negligible. It should then allow an evaluation of the Collins effect.

– $p^\uparrow p \to \pi X$, with a correct inclusion of both the Collins and Sivers effect.

Let us finally mention that we have assumed the validity of the QCD factorization theorem also at higher twist. This has been discussed in the literature: an approach similar to that discussed here [9] has been advocated, in the operator language, by Qiu and Sterman [15] who use generalized factorization theorems valid at higher twist and relate non zero single spin asymmetries in $p\,p$ collisions to the expectation value of a higher twist operator, a twist-3 parton distribution, which explicitly involves correlations between the two protons and combines quark fields with a gluonic field

strength. However, they still consider only collinear partonic configurations so that, in order to obtain non zero results, they have to take into account the contributions of higher order elementary interactions. Some more theoretical work is still necessary.

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